

An unconventional geometric phase gate with two non-identical quantum dots trapped in a photonic crystal cavity

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We propose a scheme for realizing a two-qubit controlled phase gate via an unconventional geometric phase with two non-identical and spatially separated quantum dots trapped in a photonic crystal cavity. In this system, the quantum dots simultaneously interact with a large detuned cavity mode and strong driving classical light fields. During the gate operation, the quantum dots undergo no transitions, while the cavity mode is displaced along a closed path in the phase space. In this way, the system can acquire different unconventional geometric phases conditional upon the states of the quantum dots, and a two-qubit entangling gate can be constructed. © 2011 Optical Society of America

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1. Introduction

As it is well known, a solid state implementation of cavity quantum electrodynamics (QED)-based approaches would open new opportunities for scaling the network into the practical and useful quantum information processing (QIP) systems [1]. Among the proposed schemes, the systems of self-assembled quantum dots (QDs) embedded in photonic crystal (PC) nanocavities have been a kind of very promising systems. That is not just because the strong QD-cavity interaction can be realized in these systems [2–4], but also because both QDs and PC cavities are suitable for monolithic on-chip integration.

However, there are two main challenges in this kind of systems. One is that the variation in emission frequencies of the self-assembled QDs is large [5]. It is very difficult to achieve identical QDs in experiments, so it is important to realize QIP with non-identical QDs which include both the near-resonant QDs and the nonresonant QDs. The other is that the

interaction between the QDs is difficult to control [6]. So far, there are several methods which have been used to bring the emissions of non-identical QDs into the same, such as, by using Stark shift tuning [7–9] and voltage tuning [10]. There are also several solutions which have been used to control the interaction between the QDs, for instance, coherent manipulating coupled QDs [6], and controlling the coupled QDs with Kondo effect [11]. Due to the small line widths of the QDs, cavity modes and the frequency spread of the QD ensemble, the tuning of individual QD frequencies is mainly achieved for two closely spaced QDs trapped in a PC cavity [10]. And the controlled interactions and logical gates between the QDs can be acquired in this way. Moreover, the entanglement of two QDs has been achieved in the experiment [12]. On the contrary, there are few schemes about how to achieve the controlled interactions and logical gates with the spatially separated QDs [13].

On the other hand, the quantum gates based on the dynamical phases (DPs) are sensitive to the quantum fluctuations, which are the main blockage toward a large-scale quantum computing. The ideas that adopt the geometric phase have been utilized for solving this problem [14–19]. As the geometric phase is determined only by the path area, it is insensitive to the starting state distributions, the path shape, and the passage rate to traverse the close path [15, 17]. In this aspect, geometric phases are better than the dynamical ones in realization of quantum computing. Until now, there are two kinds of geometric phase gates (GPGs). The gates containing no DP are often referred to as conventional GPGs [20]. In contrast, the gates including the DP are named unconventional GPGs [15]. Comparing with the conventional GPGs, for the DP in the unconventional GPGs being proportional to the geometric phase, they don't need to eliminate the DP. For this reason, the unconventional GPGs are better than the conventional ones. Moreover, the unconventional GPGs have been realized in the system of ions [21].

Recently, Feng *et al.* proposed a scheme to realize a quantum computation with atoms in decoherence-free subspace by using a dispersive atom-cavity interaction driven by strong classical laser fields [19]. But their proposal is based on identical qubits, and each atom is driven with four laser fields. Motivated by this work, we propose a scheme for constructing an unconventional GPG with two different QDs trapped in a single-mode cavity, and each dot is driven with two laser fields. In this scheme, the QDs undergo no transitions, while the cavity mode is displaced along a closed path in the phase space. In this way, the system can acquire different phases conditional upon different states of QDs. With the application of single-qubit operations, a controlled phase gate can be constructed in this system. During the gate operation, as the QDs undergo no transitions, the spontaneous emission of QDs can be ignored. Comparing with Ref. [19], the logical gate is extended to non-identical qubits, and the number of laser fields is decreased. The main differences between our scheme and previous ones for near-neighbor QD systems [22, 23] about the phase gates are that our

phase gate is an unconventional geometric gate, and it can be constructed with the spatially separated QDs.

The organization of this paper is as follows. In Sec.2, we introduce the theoretical model and effective Hamiltonian. In Sec.3, we review the definition of the unconventional geometric phase and present a two-qubit controlled phase gate based on the unconventional geometric phase. In Sec.4, we show the simulations and discussions of the two-qubit operation. The conclusion is given in Sec.5.

2. Theoretical Model and effective Hamiltonian

As shown Fig.1, we consider that two charged GaAs/AlGaAs QDs are non-identical and spatially separated. They are trapped in a single-mode PC cavity. Each dot has two lower states $|g\rangle = |\uparrow\rangle$, $|f\rangle = |\downarrow\rangle$ and two higher states $|e\rangle = |\uparrow\downarrow\uparrow\rangle$, $|d\rangle = |\downarrow\uparrow\downarrow\rangle$, here $(|\uparrow\rangle, |\downarrow\rangle)$ and $(|\uparrow\downarrow\uparrow\rangle, |\downarrow\uparrow\downarrow\rangle)$ denote the spin up and spin down for electron and hole, respectively. At zero magnetic field, the two lower states are twofold degenerate. The only dipole allowed transitions $|g\rangle \leftrightarrow |e\rangle$ and $|f\rangle \leftrightarrow |d\rangle$ are coupled with σ^+ and σ^- polarization lights, respectively [10, 24]. If the fields in the σ^+ polarization are applied to our system [25], the transition $|g\rangle \leftrightarrow |e\rangle$ in dot j ($= A, B$) is coupled to the cavity mode with the coupling g_j and classical laser fields with the Rabi frequencies Ω_j and Ω'_j , while $|f\rangle$ and $|d\rangle$ are not affected. The detunings for the cavity mode and classical fields are Δ_j^C , Δ_j , and $-\Delta'_j$, respectively. In our model, the quantum information is encoded in states $|g\rangle$ and $|f\rangle$. Then the Hamiltonian describing the interaction between QDs and fields is:

$$\hat{H}_I = \sum_{j=A,B} (g_j a e^{i\Delta_j^C t} + \frac{\Omega_j}{2} e^{i\Delta_j t} + \frac{\Omega'_j}{2} e^{-i\Delta'_j t}) \sigma_j^+ + H.c., \quad (1)$$

where $\sigma_j^+ = |e\rangle_j \langle g|$, a is the annihilation operator for the cavity mode.

In order to derive the effective Hamiltonian which can be used to construct a displacement operator, we use the method proposed in Refs. [19, 26, 27] and assume the following conditions: (1) $|\Omega_j| = |\Omega'_j|$; (2) $\Delta_j = \Delta'_j$; (3) the large detuning condition: $|\Delta_j|, |\Delta'_j| \gg |g_j|, |\Omega_j|, |\Omega'_j|$; (4) $\delta = \Delta_j^C - \Delta_j$; (5) $|\Omega_j| \gg |g_j|$. The first condition together with the second condition can completely cancel the Stark shifts caused by the classical light fields and related terms. Under the large detuning condition, if the QDs are initially in the ground states, since the probability for QDs absorbing photons from the light fields or being excited is negligible, the excited states will not be populated and can be adiabatically eliminated. Moreover, the third and the final conditions ensure that the terms proportional to $|g_j|^2$ and $|g_A g_B|$ can be neglected. The fourth condition can guarantee that δ is a tunable constant which is only related to detuning between the cavity field and light fields. In this situation, the effective Hamiltonian can be written as:

$$\hat{H}_{eff} = - \sum_{j=A,B} (\lambda_j a e^{i\delta t} + \lambda_j^* a^\dagger e^{-i\delta t}) |g\rangle_j \langle g|, \quad (2)$$

where $\lambda_j = \frac{\Omega_j^* g_j}{4} (\frac{1}{\Delta_j} + \frac{1}{\Delta_j^*})$. The two terms in Eq.(2) describe the coupling between the cavity mode and the classical fields induced by the virtual excited QDs.

As mentioned above, the spatially separated QDs can be addressed individually and couple independently to their corresponding classical light fields, we can assume $\lambda_j = \epsilon$ which is achievable by using laser fields with suitable Rabi frequencies. Then the effective Hamiltonian reduces to:

$$\hat{H}_{eff} = - \sum_{j=A,B} (\epsilon a e^{i\delta t} + \epsilon^* a^\dagger e^{-i\delta t}) |g\rangle_j \langle g|. \quad (3)$$

According to Refs. [15, 17–19, 21], since the effective Hamiltonian (3) is only related with the cavity mode coupling the two dots, the time evolution operator takes the form of a displacement operator, it can be used to realize two-qubit controlled phase gates.

3. Two qubits controlled phase gate

In this section, we will show how to realize the two-qubit controlled phase gate with the effective Hamiltonian (3). At first, we review the definition of the unconventional geometric phase due to a displacement operator along an arbitrary path in phase space [15, 17–19, 21]. The displacement operator is written as follows:

$$D(\alpha) = e^{\alpha a^\dagger - \alpha^* a}, \quad (4)$$

where $\alpha(\alpha^*)$ is a time-dependent parameter, a^\dagger and a are the creation and annihilation operators of the harmonic oscillator (which is the mode of the single-mode cavity in this work), respectively. According to the Baker-Campbell-Hausdorff formula, for a path consisting of N short straight sections $\Delta\alpha_m$, $m = 1, 2, 3, \dots$, the total displacement operation can be expressed as:

$$\begin{aligned} D_t &= D(\Delta\alpha_N) \dots D(\Delta\alpha_1) \\ &= \exp(i \text{Im} \{ \sum_{m=2}^N \Delta\alpha_m \sum_{k=1}^{m-1} \Delta\alpha_k^* \}) D(\sum_{m=1}^N \Delta\alpha_m). \end{aligned} \quad (5)$$

An arbitrary path γ can be followed in the limit $N \rightarrow \infty$. Thus we can get

$$D_t = e^{i\Theta} D(\int_\gamma d\alpha), \quad (6)$$

with $\Theta = \text{Im} \{ \int_\gamma \alpha^* d\alpha \}$ being the total phase which consists of both the geometric phase and the nonzero DP. Since the nonzero DP is proportional to the geometric phase, the total phase is an unconventional phase [15].

When the path is closed, we can obtain

$$D_t = D(0)e^{i\Theta} = e^{i\Theta}, \quad (7)$$

$$\Theta = \text{Im}\{\oint \alpha^* d\alpha\}. \quad (8)$$

In this case, the phase Θ is determined by the area of a loop in the phase space and independent of the quantized state of the harmonic oscillator [17].

According to the definition of the unconventional geometric phase, in the infinitesimal interval $[t, t + dt]$, the system governed by the effective Hamiltonian Eq.(3) will evolve as follows:

$$\begin{cases} |ff\rangle|\theta_{ff}(t)\rangle \rightarrow |ff\rangle|\theta_{ff}(t)\rangle, \\ |fg\rangle|\theta_{fg}(t)\rangle \rightarrow e^{-iH_{eff}t}|fg\rangle|\theta_{fg}(t)\rangle = D(d\alpha_{fg})|fg\rangle|\theta_{fg}(t)\rangle, \\ |gf\rangle|\theta_{gf}(t)\rangle \rightarrow D(d\alpha_{gf})|gf\rangle|\theta_{gf}(t)\rangle, \\ |gg\rangle|\theta_{gg}(t)\rangle \rightarrow D(d\alpha_{gg})|gg\rangle|\theta_{gg}(t)\rangle, \end{cases} \quad (9)$$

where $d\alpha_{fg} = i\epsilon^*e^{-i\delta t}dt$, $d\alpha_{gf} = d\alpha_{fg}$, $d\alpha_{gg} = 2d\alpha_{fg}$, and $|\theta_{u,v}(t)\rangle (u, v = g, f)$ denotes the state of the cavity mode, which is determined by the qubit state $|u_A\rangle|v_B\rangle$ at the time t .

If the cavity field is initially in the vacuum state $|0\rangle$, after an interaction time t , the evolution of the system takes the form of

$$\begin{cases} |ff\rangle|0\rangle \rightarrow |ff\rangle|0\rangle, \\ |fg\rangle|0\rangle \rightarrow e^{i\phi_{fg}}D(\alpha_{fg})|fg\rangle|0\rangle, \\ |gf\rangle|0\rangle \rightarrow e^{i\phi_{gf}}D(\alpha_{gf})|gf\rangle|0\rangle, \\ |gg\rangle|0\rangle \rightarrow e^{i\phi_{gg}}D(\alpha_{gg})|gg\rangle|0\rangle, \end{cases} \quad (10)$$

with

$$\begin{cases} \alpha_{fg} = i \int_0^t \epsilon^* e^{-i\delta t} dt = -\frac{\epsilon^*}{\delta}(e^{-i\delta t} - 1), \\ \alpha_{gf} = i \int_0^t \epsilon^* e^{-i\delta t} dt = -\frac{\epsilon^*}{\delta}(e^{-i\delta t} - 1), \\ \alpha_{gg} = i \int_0^t 2\epsilon^* e^{-i\delta t} dt = \alpha_{gf} + \alpha_{fg}, \end{cases} \quad (11)$$

and

$$\begin{cases} \phi_{fg} = \text{Im}(\int \alpha_{fg}^* d\alpha_{fg}) = -\frac{|\epsilon|^2}{\delta}(t - \frac{\sin(\delta t)}{\delta}), \\ \phi_{gf} = \text{Im}(\int \alpha_{gf}^* d\alpha_{gf}) = -\frac{|\epsilon|^2}{\delta}(t - \frac{\sin(\delta t)}{\delta}), \\ \phi_{gg} = \text{Im}(\int \alpha_{gg}^* d\alpha_{gg}) = \phi_{gf} + \phi_{fg} + \theta_{gg} \end{cases} \quad (12)$$

$$\begin{aligned} \theta_{gg} &= \text{Im}(\int \alpha_{gf}^* d\alpha_{fg} + \int \alpha_{fg}^* d\alpha_{gf}) \\ &= \text{Im}(-2 \int_0^t \frac{|\epsilon|^2}{\delta}(1 - e^{-i\delta t})i dt) \\ &= -\frac{2|\epsilon|^2}{\delta}(t - \frac{\sin(\delta t)}{\delta}). \end{aligned} \quad (13)$$

Eq.(11) shows, under the condition $t = 2l\pi/\delta$, for $l = 1, 2, 3, \dots$, the displacement parameter $d\alpha_{uv}$ for state $|u_A\rangle|v_B\rangle$ moves along a closed path and returns to the original point in the

phase space of the coherent state $|\alpha_{uv}\rangle$. And the system can acquire the different unconventional geometric phases (12) conditional upon the different states of QDs in this way. Within a definite period of time, e.g. from $t = 2l\pi/\delta$ to $t = 2(l+1)\pi/\delta$, the state for the cavity mode evolves from a vacuum state to a coherence state at first, and then evolves to the vacuum state again. Moreover, since Eq.(3) is only dependent on the detuning δ which is not related to the energy-level of QDs, the two-qubit operation (10) can be realized with the system constituted of both non-resonant and near-resonant QDs.

Owing to the detuning δ , after an interaction time $t = 2l\pi/\delta$, the evolution of the wavefunction for the system (10) can be summarized by [21]:

$$\begin{cases} |ff\rangle|0\rangle \rightarrow |ff\rangle|0\rangle, \\ |fg\rangle|0\rangle \rightarrow e^{-\Phi}|fg\rangle|0\rangle, \\ |gf\rangle|0\rangle \rightarrow e^{-\Phi}|gf\rangle|0\rangle, \\ |gg\rangle|0\rangle \rightarrow e^{-4\Phi}|gg\rangle|0\rangle. \end{cases} \quad (14)$$

where $\Phi = 2l\pi|\epsilon|^2/\delta^2$. This two-qubit operation is induced by the cavity mode which couples the two dots. In this situation, since the displacement parameters satisfy $d\alpha_{gg} = 2d\alpha_{gf} = 2d\alpha_{fg}$, according to Eqs.(11) and (12), the relationship for the amplifies of the coherent states is $\alpha_{gg} = 2\alpha_{fg} = 2\alpha_{gf}$, and the closed area for ϕ_{gg} is four times of the one for $\phi_{fg} = \phi_{gf}$. So unconventional geometric phases satisfy $\phi_{gg} = 4\phi_{fg} = 4\phi_{gf}$.

After the application of the single-qubit operations $|g\rangle_j \rightarrow e^{\Phi}|g\rangle_j$ [28], two-qubit operation (14) transforms into

$$\begin{cases} |ff\rangle|0\rangle \rightarrow |ff\rangle|0\rangle, \\ |fg\rangle|0\rangle \rightarrow |fg\rangle|0\rangle, \\ |gf\rangle|0\rangle \rightarrow |gf\rangle|0\rangle, \\ |gg\rangle|0\rangle \rightarrow e^{-2\Phi}|gg\rangle|0\rangle. \end{cases} \quad (15)$$

The above equation represents an unconventional geometric phase gate. In this quantum phase gate operation, if and only if both qubits are in the state $|g\rangle$, there will be an additional phase -2Φ in the system.

4. Simulations and Discussions

Finally, we present some numerical simulations to demonstrate our proposal can be realized in current experimental setups and it can robust against prevailing uncertainties and fluctuations of parameters to some extent. For the sake of convenience, we take the two-qubit operation (14) with $\Phi = \pi/2$ as an example to discuss the realization. Under the condition of large detuning, the influences of spontaneous emission from the excited states of QDs can be ignored, and the main decoherence effect in this system is due to cavity decay. Then the

master equation can be written as follows:

$$\dot{\rho} = -i[H_I, \rho] + \frac{\gamma}{2}(2a\rho a^\dagger - a^\dagger a\rho - \rho a^\dagger a), \quad (16)$$

where ρ is the reduced density operator of the system, γ is the cavity decay rate.

The fidelity of the two-qubit operation can be expressed as:

$$F = \langle \Psi | \rho(T) | \Psi \rangle, \quad (17)$$

where $T = \pi\delta/(2|\epsilon|^2)$ is the shortest two-qubit operation time, $|\Psi\rangle$ is a target state which is dependent on the initial state. For instance, if the initial state is in the state $|\Psi_{in}\rangle = \frac{(x|ff\rangle + y|gf\rangle + z|fg\rangle + w|gg\rangle)}{\sqrt{x^2 + y^2 + z^2 + w^2}}$ with random coefficients $\{x, y, z, w\}$, the corresponding target state takes the form of $|\Psi\rangle = \frac{(x|ff\rangle - iy|gf\rangle - iz|fg\rangle + w|gg\rangle)}{\sqrt{x^2 + y^2 + z^2 + w^2}}$. The numerical calculations for the fidelities of the two-qubit operation versus the cavity decay rate are given in Fig. 2, and these fidelities are the average over fidelities for 500 different initial states $|\Psi_{in}\rangle$. In doing so, some experimental parameters in Refs. [24, 29, 30] are referred.

Fig. 2 shows the following: Firstly, with the increase of γ/γ_0 , here $\gamma_0 = (5ns)^{-1}$ is the cavity decay rate that has been achieved in the experiment [29, 30], the fidelity for the two-qubit operation decreases. It means that the cavity decay affects the fidelity of the two-qubit operation largely [19]. The reason for this is that the state of cavity mode evolves between the vacuum state and the coherent state. It is worth pointing out that the decay of coherent state depends on the mean photon number of coherent state and the cavity decay. On the one hand, when the mean photon number is definite, the decay of coherent state increases with increasing the cavity decay. On the other hand, when the cavity decay is definite, the decay of coherent state increases with increasing the mean photon number of coherent state. Secondly, with the increase of δ , the fidelity of the two-qubit operation increases for the decrease of the mean photon number of coherent state. Thirdly, the fidelity for $\delta = 0.25g_A$ ($\delta = 2g_A$) is about 99.98%(99.95%) when the cavity decay rate is $\gamma = \gamma_0$, and it decreases to about 99.96%(99.88%) when $\gamma = 2\gamma_0$. Moreover, as accumulated unconventional geometric phase for one loop would be small, our system has to take multi-loops. For this reason, our scheme needs a good cavity, which can prevent the photons leaking from the cavity and ensure the higher fidelity. In addition, according to the parameters in Fig. 2, the two-qubit operation time for $\delta = 0.25g_A$ ($\delta = 2g_A$) is about 1.7ns(13.5ns) which is much smaller than the effective decay time of cavity $\gamma/(\frac{|\epsilon|^2}{\delta^2}) \sim 120ns(4000ns)$. Therefore, it is possible to realize our scheme in the experiment.

Besides the effect of the decoherence, there are some other factors that affect the fidelity of the two-qubit operation, the prevailing factors in our scheme are the uncertainties and fluctuations of parameters, for example, they might add unwanted phases to the two-qubit operation. According to Ref. [31], we can calculate the fidelity of the two-qubit operation

versus the parameter fluctuation as shown in Fig.3. For the sake of simplicity, we assume $\zeta = \Delta k_j/k_j > 0$ with Δk_j being the corresponding parameter fluctuation. In this case, both QDs undergo the same parameter fluctuations when ζ is definite. Fig.3 shows that the variation trends for all curves of the fidelity are the same. In the absence of parameter fluctuations, the fidelities of two-qubit operation can be as high as $F = 99.95\%$. When the parameter fluctuations are included, fidelities decrease with the increase of parameter fluctuations. The fidelity for $\zeta = 0.02$ is about 99.62% , and it decreases to about 98.81% when $\zeta = 0.04$. Therefore, to some extent our scheme can also resist against errors due to the parameter fluctuations. Note that, each curve in Fig.3 is only corresponding to one kind of parameter fluctuation.

5. Conclusion

In conclusion, we have shown that in a single-mode PC cavity, two non-identical and spatially separated QDs driven by the classical light fields can be used to realize the two-qubit controlled phase gate. During the gate operation, the QDs remain in their ground states, while the cavity mode is displaced along a circle in the phase space, and the system can acquire different unconventional geometric phases conditional upon the states of QDs. After the application of the single-qubit operations, the controlled phase gate based on the unconventional geometric phase can be constructed. The distinct advantages of the proposed scheme are the follows: firstly, as this controlled phase gate is based on the unconventional geometric phase, it can resist against the parameter uncertainties and fluctuations to some extent [15, 21]; secondly, as the QDs are non-identical, it is more practical in the experiment; thirdly, as the quantum information is encoded in the two ground states, this gate is insensitive to the spontaneous emission of the QDs [17]. Finally, we have to point out the shortcoming in our scheme is that the single qubit operations cannot be realized with our model directly.

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List of Figure Captions

Fig. 1. (Color online) (a) Schematic diagram of the system. (b) the configuration of the QDs level structure and relevant transitions. The states $|g\rangle$ and $|f\rangle$ correspond to two lower levels, while $|e\rangle$ and $|d\rangle$ are two higher levels. The transition $|g\rangle \leftrightarrow |e\rangle$ for each dot is driven by the cavity field and the classical pulses with the detunings Δ_j^C , Δ_j and Δ_j' , respectively. g_j represents the coupling rate of the QDs to cavity mode, Ω_j and Ω_j' are the Rabi frequency of the classical pulses.

Fig. 2. (Color online) Numerical simulation of the fidelity of the two-qubit operation (14) versus the cavity decay, with the parameters $g_A = 0.10meV$, $g_B = 0.08meV$, $\Omega_A = 10meV$, $\Omega_B = 13.75meV$, $\gamma_0 = (5ns)^{-1}$. The detunings of blue line are given by $\Delta_A = 200.00meV$, $\Delta_B = 220.00meV$, and the detunings of green line are given by $\Delta_A = 200.09meV$, $\Delta_B = 220.09meV$, respectively.

Fig. 3. (Color online) Numerical simulation of the fidelity of the two-qubit operation (14) versus the parameter fluctuations. The decay rate for cavity is given by $\gamma = \gamma_0$. All other parameters are the same for the blue line in Fig. 2.

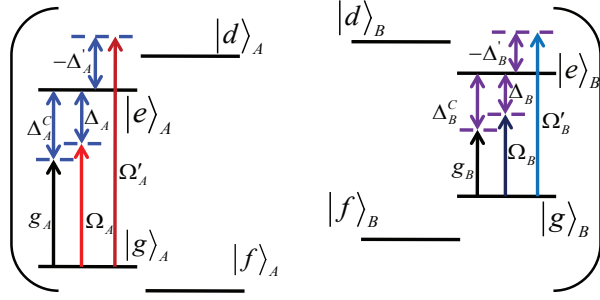


Fig. 1. (Color online) (a) Schematic diagram of the system. (b) The configuration of the level structure and relevant transitions for dot j ($= A, B$). The states $|g\rangle_j$ and $|f\rangle_j$ correspond to two lower levels, while $|e\rangle_j$ and $|d\rangle_j$ are two higher levels. The transition $|g\rangle_j \leftrightarrow |e\rangle_j$ is coupled to the cavity mode with g_j and classical laser fields with the Rabi frequencies Ω_j and Ω'_j . The detunings for the cavity mode and classical fields are Δ_j^C , Δ_j , and $-\Delta'_j$, respectively. All the light fields are in the σ^+ polarization.

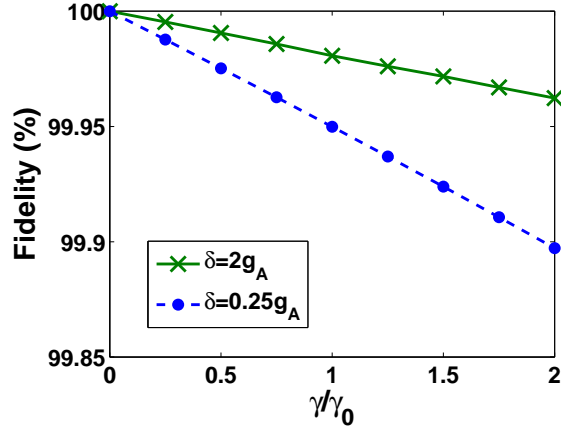


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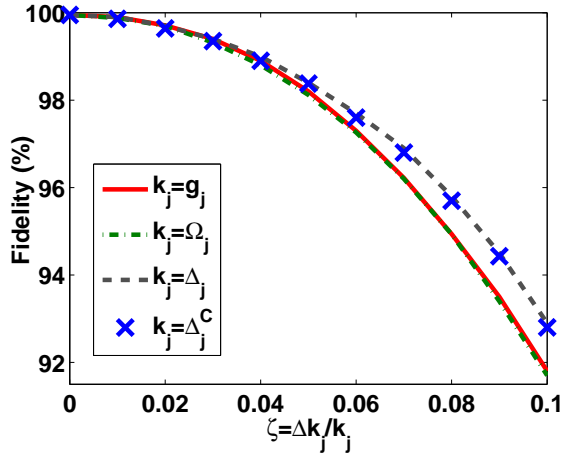


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